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Classical theory of canonical QCD on a space-like hypersurface

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The canonical formalism in classical theory of QCD is constructed on a space-like hypersurface. The Poisson bracket on the space-like hypersurface is defined and it plays an important role to describe every algebraic relation in the canonical formalism into Lorentz covariant form. Surface integrals are introduced as alternatives of field equations for quarks, gluons, and Faddeev-Popov ghosts. It is shown that deformations of the space-like hypersurface for surface integrals are generated by the interaction term of QCD Hamiltonian density. By converting the Poisson bracket on the space-like hypersurface to four-dimensional commutator, we can pass over to QCD in the Heisenberg picture without spoiling the explicit Lorentz covariance.

One of the problems in the ordinary canonical method of quantization in quantum chromodynamics (QCD) is in the classical procedure blinding the Lorentz covariance. The definition of the canonical conjugate momentum depends on the frame of reference. The Hamiltonian is linked to a definite choice of the time. Further, the commutation relation of two field variables is defined at the same time. From the relativistic point of view, the canonical formalism on the basis of a particular frame of reference spoils the explicit Lorentz covariance in the classical description of QCD far from field quantization. For this reason the canonical method of quantization in QCD have been regarded to be inferior to path-integral method. In the path-integral method Lorentz covariance is manifest from the outset. This is an impressive method to construct the S-matrix in QCD with covariant manner. Although the canonical method is superior to the path-integral method in presenting clearly a particle interpretation to quarks, gluons, and Faddeev-Popov ghosts, the lack of explicit Lorentz covariance turns our interest to the path-integral method. This paper presents a canonical formalism in the classical theory of QCD with coherently Lorentz covariant form. The main subject is to construct the canonical formalism on a space-like hypersurface. It is shown that the canonical method of quatization in QCD should be applied on the space-like hypersurface by converting classical fields to quantized field operators in the Heisenberg picture. The generalization of our formulation will be discussed in separate papers.

The total QCD Lagrangian density with $\xi = 1$ gauge is taken as

$$\begin{aligned}\mathcal{L}[x] = {} & \bar{\psi}(x) \left(i\gamma^\mu \partial_\mu - m + g\gamma^\mu \lambda^i W^i_\mu(x)\right) \psi(x) - \frac{1}{4} G^i_{\mu\nu}(x) G^{i\mu\nu}(x) - \frac{1}{2}\left(\partial^\mu W^i_\mu(x)\right)^2 \\ & + \frac{1}{2}\partial_\mu \left[W^i_\nu(x)\left(g^{\mu\nu}\partial_\rho W^{i\rho}(x) - \partial_\nu W^{i\mu}(x)\right)\right] + \left(\partial^\mu \bar{c}^a(x)\right) D^{ab}_\mu c^b(x), \qquad (1)\end{aligned}$$

with

$$\begin{aligned} G^i_{\mu\nu}(x) &= \partial_\mu W^i_\nu(x) - \partial_\nu W^i_\mu(x) + g f^{ijk} W^j_\mu(x) W^k_\nu(x), \\ D^{ab}_\mu &= \delta^{ab}\partial_\mu + g f^{acb} W^c_\mu(x), \end{aligned}$$

where ψ is the $SU_c(3)$ triplet quark, W^i_μ are gluon fields, c^a and $\bar{c}^a$ are Faddeev-Popov ghosts and anti-ghosts, λ^i are the Gell-Mann matrices, and f^{ijk} are the structure constants of $SU_c(3)$. With the Lagrangian density (1), the principle of the least action gives field equations

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$$\left(i\gamma^\mu\partial_\mu - m + g\gamma^\mu\lambda^i W^i_\mu(x)\right)\psi(x) = 0, \tag{2}$$

$$\partial_\mu\left[\partial^\mu W^{i\rho}(x) + gf^{ijk}W^{j\mu}(x)W^{k\rho}(x)\right] - gf^{nik}W^k_\nu(x)G^{n\rho\nu}(x)$$
$$+gf^{iab}[\partial^\rho\bar{c}^a(x)]c^b(x) - g\bar{\psi}(x)\gamma^\rho\lambda^i\psi(x) = 0, \tag{3}$$

$$\partial^\mu D^{ab}_\mu c^b(x) = 0, \tag{4}$$

$$D^{ab\mu}\partial_\mu\bar{c}^b(x) = 0. \tag{5}$$

In QCD all field variables appeared in the Lagrangian density are field operators. However, we classicize the gluon field operators into Grassmann even wave functions, and the quark field and Faddeev-Popov ghost operators into Grassmann odd wave functions to look into only the classical aspect of QCD.

The canonical formalism needs the canonical variables. We will adopt every classical field and its conjugate momentum on a space-like hypersurface as the canonical variables. Let us choose conjugate momenta on the space-like hypersurface σ as

$$\Pi_\psi(x) = n_\mu(x)\frac{\partial^R\mathcal{L}[x]}{\partial\left[\partial_\mu\psi(x)\right]} = in_\mu(x)\bar{\psi}(x)\gamma^\mu, \tag{6}$$

$$\Pi^{i\nu}_W(x) = n_\mu(x)\frac{\partial^R\mathcal{L}[x]}{\partial\left[\partial_\mu W^i_\nu(x)\right]} = -n^\mu(x)\partial_\mu W^{i\nu}(x) - gf^{ijk}(n^\mu(x)W^j_\mu(x))W^{k\nu}(x), \tag{7}$$

$$\Pi^a_c(x) = n_\mu(x)\frac{\partial^R\mathcal{L}[x]}{\partial\left[\partial_\mu c^a(x)\right]} = n^\mu(x)\partial_\mu\bar{c}^a(x), \tag{8}$$

$$\Pi^a_{\bar{c}}(x) = n_\mu(x)\frac{\partial^R\mathcal{L}[x]}{\partial\left[\partial_\mu\bar{c}^a(x)\right]} = -n^\mu(x)\partial_\mu c^a(x) - gf^{acb}\left(n^\mu(x)W^c_\mu(x)\right)c^b(x), \tag{9}$$

where $n^\mu(x)$ is a unit normal ($n^\mu(x)n_\mu(x) = 1$) at x^μ to the surface passing through x^μ. The superscript " R " means that we have adopted the right-differentiation convention. If we make the space-like hypersurface to be global flat surface, the conjugate momenta (6)–(9) reduce to the ordinary conjugate momenta. Thus our conjugate momenta are natural extension of the conventional ones defined at a particular time.

The total QCD Hamiltonian density is given by the Legendre transformation on σ:

$$\begin{aligned}\mathcal{H}[x] &= \Pi_\psi(x)\partial_n\psi(x) + \Pi^{i\nu}_W(x)\partial_n W^i_\nu(x) + \Pi^a_c(x)\partial_n c^a(x) + \Pi^a_{\bar{c}}(x)\partial_n\bar{c}^a(x) - \mathcal{L}[x] \\ &= \mathcal{H}_0[x] + \mathcal{H}_{\text{int}}[x],\end{aligned} \tag{10}$$

with

$$\begin{aligned}\mathcal{H}_0[x] = &- \Pi_\psi(x)\left(n\cdot\gamma\right)\gamma_\mu\partial^\mu_t\psi(x) - i\Pi_\psi(x)\left(n\cdot\gamma\right)\gamma^\nu\psi(x) \\ &- \frac{1}{2}\Pi_{W}{}^i_\mu(x)\Pi_W{}^{i\mu}(x) + \frac{1}{2}\left(\partial_{t\mu}W^{i\nu}(x)\right)\left(\partial^\mu_t W^i_\nu(x)\right) \\ &+ \Pi^a_{\bar{c}}(x)\Pi^a_c(x) - \left(\partial_{t\mu}\bar{c}^a(x)\right)\left(\partial^\mu_t\bar{c}^a(x)\right),\end{aligned}$$

$$\begin{aligned}\mathcal{H}_{\text{int}}[x] = \quad & ig\Pi_\psi(x)(n\cdot\gamma)\gamma^\mu\lambda^i\psi(x)W^i_\mu(x) \\ &- g\left[f^{ilm}(n\cdot W^l)\Pi^i_{W\nu}(x)W^{m\nu}(x) - \frac{1}{2}f^{ijk}W^j_\mu(x)W^k_\nu(x)\left(\partial^\nu_t W^{i\nu}(x) - \partial^\nu_t W^{i\mu}(x)\right)\right] \\ &+ \frac{g^2}{2}f^{ijk}f^{ilm}\left[(n\cdot W^k)(n\cdot W^l)W^j_\mu(x)W^{m\mu}(x) + \frac{1}{2}W^j_\mu(x)W^k_\nu(x)W^{l\mu}(x)W^{m\nu}(x)\right] \\ &- gf^{acb}\Pi^a_c(x)\left(n\cdot W^c\right)c^b(x) - gf^{acb}\left(\partial_{t\nu}\bar{c}^a(x)\right)W^{c\nu}(x)c^b(x),\end{aligned}$$

where

$$\partial_\mu = n_\mu(x)\partial_n + \partial_{t\mu},$$

$$\partial_n = n_\mu(x)\partial^\mu, \quad \partial_{t\mu} = (g_{\mu\nu} - n_\mu(x)n_\nu(x))\partial^\nu.$$

The differential operators ∂_n and $\partial_{t\mu}$ are the directional derivatives in the direction of normal and tangent to σ, respectively. With the help of the total canonical energy-momentum tensor $T_{\mu\nu}[x]$, we can express (10) in the form

$$\mathcal{H}[x] = n^\mu(x)n^\nu(x)T_{\mu\nu}[x].$$

Thus we can identify $\mathcal{H}_{0[x]}$ as the kinetic term of the Hamiltonian density, and $\mathcal{H}_{\mathrm{int}[x]}$ as the interaction term of the Hamiltonian density. The surface integral of $\mathcal{H}$ on σ should be regarded as the total QCD Hamiltonian:

$$H = \int_\sigma d\Sigma^\mu{}_{(x)} n^\nu{}_{(x)} T_{\mu\nu[x]} = \int_\sigma d\Sigma \mathcal{H},$$

where $d\Sigma^\mu{}_{(x)} = n^\mu{}_{(x)} d\Sigma$. We get Hamilton's equations by taking the variation of H on σ:

$$\frac{\partial^R \mathcal{H}}{\partial \Pi_\psi} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} \Pi_\psi\right]} + \partial_n \psi = 0, \qquad \frac{\partial^R \mathcal{H}}{\partial \psi} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} \psi\right]} + \partial_n \Pi_\psi = 0, \tag{11}$$

$$\frac{\partial^R \mathcal{H}}{\partial W^i_\lambda} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} W^i_\lambda\right]} + \partial_n \Pi^{i\lambda}_W = 0, \; \frac{\partial^R \mathcal{H}}{\partial \Pi^i_{W\lambda}} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} \Pi^i_{W\lambda}\right]} - \partial_n W^{i\lambda} = 0, \tag{12}$$

$$\frac{\partial^R \mathcal{H}}{\partial \Pi^a_c} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} \Pi^a_c\right]} + \partial_n c^a = 0, \qquad \frac{\partial^R \mathcal{H}}{\partial c^a} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} c^a\right]} + \partial_n \Pi^a_c = 0, \tag{13}$$

$$\frac{\partial^R \mathcal{H}}{\partial \Pi^a_{\bar{c}}} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} \Pi^a_{\bar{c}}\right]} + \partial_n \bar{c}^a = 0, \qquad \frac{\partial^R \mathcal{H}}{\partial \bar{c}^a} - \partial_{t\mu} \frac{\partial^R \mathcal{H}}{\partial \left[\partial_{t\mu} \bar{c}^a\right]} + \partial_n \Pi^a_{\bar{c}} = 0. \tag{14}$$

They give field equations (2)–(5) (excepting a half set of equations for gluons leading to a trivial identity $0 = 0$), so Hamilton's equations (11)–(14) are equivalent to the Euler-Lagrange equations.

Hamilton's equations (11)–(14) and the Klein-Gordon equation associated with the invariant delta function ($(\Box + m^2)\Delta(x) = 0$; $\Box D(x) = 0$) can be combined to give the following surface integrals

$$\psi(x;\sigma) = -\int_\sigma d\Sigma (i\gamma\cdot\partial + m)\Delta(x - x'')\gamma_0 \Pi^\dagger_\psi(x''), \tag{15}$$

$$\bar{\psi}(x;\sigma) = -\int_\sigma d\Sigma \Pi_\psi(x'')(-i\gamma\cdot\partial + m)\Delta(x - x''), \tag{16}$$

$$W^{i\lambda}(x;\sigma) = \int_\sigma d\Sigma [\partial''_n D(x - x'') \cdot W^{i\lambda}(x'') + D(x - x'')\Pi^{i\lambda}_W(x'')], \tag{17}$$

$$c^a(x;\sigma) = \int_\sigma d\Sigma [\partial''_n D(x - x'') \cdot c^a(x'') + D(x - x'')\Pi^a_{\bar{c}}(x'')], \tag{18}$$

$$\bar{c}^a(x;\sigma) = \int_\sigma d\Sigma [\partial''_n D(x - x'') \cdot \bar{c}^a(x'') - D(x - x'')\Pi^a_c(x'')]. \tag{19}$$

Note that the world-point x is not necessarily on the surface σ, and space-time coordinates on σ are represented by x''. We find that the integrals (15)–(19) satisfy free field equations:

$$\begin{aligned} (i\gamma\cdot\partial - m)\ \psi(x;\sigma) &= 0, \\ \Box\ W^{i\lambda}(x;\sigma) &= 0, \\ \Box\ c^a(x;\sigma) &= 0, \\ \Box\ \bar{c}^a(x;\sigma) &= 0. \end{aligned}$$

Thus the surface integrals (15)–(19) are equivalent to asymptotic fields.

The bubble differentiation [1–3] (a differentiation with respect to σ) of the integrals (15)–(19) around a world-point x' reads

$$\frac{\delta\psi(x;\sigma)}{\delta\sigma_{(x')}} = -g(i\gamma\cdot\partial + m)\Delta(x - x')\gamma^\mu\lambda^i\psi_{(x')}W^i_{\mu(x')}, \tag{20}$$

$$\begin{aligned} \frac{\delta W^{i\lambda}(x;\sigma)}{\delta\sigma_{(x')}} = &-\partial'^\mu D(x - x')\cdot g f^{ijk} W^j_{\mu(x')} W^{k\lambda}{}_{(x')} + D(x - x')\cdot g f^{ijk} G^{j\lambda\nu}{}_{(x')} W^k_{\nu(x')} \\ &+ D(x - x')\cdot g f^{icb}[\partial'^\lambda \bar{c}^c{}_{(x')}] c^b{}_{(x')} - D(x - x')\cdot ig\bar{\psi}_{(x')}\gamma^\lambda\lambda^i\psi_{(x')}, \end{aligned} \tag{21}$$

$$\frac{\delta c^a(x;\sigma)}{\delta\sigma_{(x')}} = -g f^{acb}\partial'^\mu D(x - x')\cdot W^c_{\mu(x')} c^b{}_{(x')}, \tag{22}$$

$$\frac{\delta\bar{c}^a(x;\sigma)}{\delta\sigma_{(x')}} = g f^{acb} D(x - x') W^c_{\mu(x')}\partial'^\mu \bar{c}^b{}_{(x')}. \tag{23}$$

The right-hand side of equations (20)–(23) includes only the interaction term, so the integrals (15)–(19) are not conserved because of the interaction between fields. However, equations (20)–(23) are invariant under the canonical transformations. The invariance is easily verified to represent equations (20)–(23) in terms of Hamilton's equations and the Klein-Gordon equation of the invariant delta function.

If we impose the Lorentz condition $\partial''_\rho W^{i\rho}(x'') = 0$ on σ as a boundary condition, the supplementary condition of the $W^{i\rho}(x;\sigma)$ becomes

$$\begin{aligned}\partial_\rho W^{i\rho}(x;\sigma) - \int_\sigma d\Sigma_{\rho(x'')}[\partial''^\mu D(x-x'') \cdot g f^{ijk} W^j_\mu{}_{(x'')} W^{k\rho}{}_{(x'')} \\ - D(x-x'') \cdot g f^{ijk} G^{j\rho\nu}{}_{(x'')} W^k_\nu{}_{(x'')} \\ - D(x-x'') \cdot g f^{icb} [\partial'^\lambda \bar{c}^c{}_{(x'')}] c^b{}_{(x'')} \\ + D(x-x'') \cdot ig \bar{\psi}_{(x'')} \gamma^\rho \lambda^i \psi_{(x'')}] \approx 0.\end{aligned} \tag{24}$$

The weak equality symbol $\approx$ is to emphasize that the left-hand side is numerically restricted to be zero but does not identically vanish in the phase space.

Now, we define a Poisson bracket on the space-like hypersurface σ. Any two Lorentz covariant quantities $F[\sigma]$ and $G[\sigma]$ being the function of canonical field variables and momenta on σ have a Poisson bracket which we shall denote by $[F[\sigma], G[\sigma]]_c$, defined by

$$[F[\sigma], G[\sigma]]_c = \int_\sigma d\Sigma \left(\frac{\tilde{\delta} F[\sigma]}{\tilde{\delta}\phi_{A(z)}} \frac{\tilde{\delta} G[\sigma]}{\tilde{\delta}\Pi_{A(z)}} - (-1)^{|A|} \frac{\tilde{\delta} F[\sigma]}{\tilde{\delta}\Pi_{A(z)}} \frac{\tilde{\delta} G[\sigma]}{\tilde{\delta}\phi_{A(z)}} \right). \tag{25}$$

The $|A|$ is the number of factor associated with the classical field ϕ_A: take 0 for Grassmann even function, and 1 for Grassmann odd function. If $F[\sigma]$ is given by

$$\begin{aligned}F[\sigma] &= \int_\sigma d\Sigma^\nu{}_{(x)}\ f_\nu(\phi_{A(x)}, \partial_{\mu t}\phi_A(x), \Pi_{A(x)}, \partial_{\mu t}\Pi_{A(x)}) \\ &= \int_\sigma d\Sigma\ \mathcal{F}(\phi_A, \partial_{\mu t}\phi_A, \Pi_A, \partial_{\mu t}\Pi_A)\end{aligned}$$

with $d\Sigma^\nu f_\nu = d\Sigma(n \cdot f) \equiv d\Sigma\mathcal{F}$, the functional derivative on σ reads

$$\frac{\tilde{\delta} F[\sigma]}{\tilde{\delta}\phi_{A(z)}} = \frac{\partial^R \mathcal{F}}{\partial\phi_{A(z)}} - \partial_{t\mu} \frac{\partial^R \mathcal{F}}{\partial[\partial_{t\mu}\phi_{A(z)}]},$$

$$\frac{\tilde{\delta} F[\sigma]}{\tilde{\delta}\Pi_{A(z)}} = \frac{\partial^R \mathcal{F}}{\partial\Pi_{A(z)}} - \partial_{t\mu} \frac{\partial^R \mathcal{F}}{\partial[\partial_{t\mu}\Pi_{A(z)}]}.$$

Our Poisson bracket (25) is a generalization of the ordinary equal-time Poisson bracket to be Lorentz covariant. All algebraic relations of the ordinary Poisson bracket hold for the Poisson bracket on σ. Here we will call the Poisson bracket on σ the four-dimensional Poisson bracket.

Substituting (15)–(18) into (25), we get the four-dimensional Poisson bracket relations of $\psi(x;\sigma), W^i_\mu(x;\sigma), c^a(x;\sigma)$, and $\bar{c}^a(x;\sigma)$:

$$\left[\psi_\alpha(x;\sigma), \bar{\psi}_\beta(y;\sigma)\right]_c = (i\gamma \cdot \partial + m)_{\alpha\beta} \Delta(x-y), \tag{26}$$

$$\left[W^i_\mu(x;\sigma), W^j_\nu(y;\sigma)\right]_c = -\delta^{ij} g_{\mu\nu} D(x-y), \tag{27}$$

$$\left[c^a(x;\sigma), \bar{c}^b(y;\sigma)\right]_c = \delta^{ab} D(x-y). \tag{28}$$

If there are no interaction, $\psi(x;\sigma), W^i_\mu(x;\sigma), c^a(x;\sigma)$, and $\bar{c}^a(x;\sigma)$ are conserved and do not depend on σ. It follows that

$$\begin{aligned}\left[\psi_\alpha(x), \bar{\psi}_\beta(y)\right]_c &= (i\gamma \cdot \partial + m)_{\alpha\beta} \Delta(x-y), \\ \left[W^i_\mu(x), W^j_\nu(y)\right]_c &= -\delta^{ij} g_{\mu\nu} D(x-y), \\ \left[c^a(x), \bar{c}^b(y)\right]_c &= \delta^{ab} D(x-y).\end{aligned}$$

These four-dimensional Poisson bracket relations correspond to the four-dimensional commutation relations for free fields in QCD.

Equations (20)–(23) are also expressible in terms of the four-dimensional Poisson bracket:

$$\frac{\delta\psi(x;\sigma)}{\delta\sigma(x')} = -\left[\psi(x;\sigma), \mathcal{H}_{\mathrm{int}\,[x']}\right]_c , \tag{29}$$

$$\frac{\delta W^{i\lambda}(x;\sigma)}{\delta\sigma(x')} = \left[W^{i\lambda}(x;\sigma), \mathcal{H}_{\mathrm{int}\,[x']}\right]_c , \tag{30}$$

$$\frac{\delta c^a(x;\sigma)}{\delta\sigma(x')} = -\left[c^a(x;\sigma), \mathcal{H}_{\mathrm{int}\,[x']}\right]_c , \tag{31}$$

$$\frac{\delta \bar{c}^a(x;\sigma)}{\delta\sigma(x')} = -\left[\bar{c}^a(x;\sigma), \mathcal{H}_{\mathrm{int}\,[x']}\right]_c . \tag{32}$$

These equations show that the displacement of $\psi(x;\sigma), W^i_\mu(x;\sigma), c^a(x;\sigma)$ and $\bar{c}^a(x;\sigma)$ under the deformation of σ is equal to that of the surface integrals under the canonical transformation generated by the interaction Hamiltonian density $\mathcal{H}_{\mathrm{int}}$ at a world-point x' on σ.

To pass over to QCD in the Heisenberg picture, we shall make all canonical variables into operators and convert all the four-dimensional Poisson bracket relations to four-dimensional commutation relations:

$$\text{(four-dimensional Poisson bracket)} \rightarrow (i\hbar)^{-1}\text{(four-dimensional commutator)}.$$

In this way we can pass over to QCD without spoiling the Lorentz covariance. A remarkable feature of our prescription in quantization is that equations (29)–(32) are converted to the Yang–Feldman equations [4]. According to the procedure given by Yang and Feldman, the S-matrix of the Heisenberg picture in QCD must be identified as

$$\mathbf{S} = \cdots \left[1 - i\int_{\sigma_{-1}}^{\sigma_0} \widehat{\mathcal{H}}_{\mathrm{int}\,[x']}\; d^4x'\right]\left[1 - i\int_{\sigma_0}^{\sigma_1} \widehat{\mathcal{H}}_{\mathrm{int}\,[x']}\; d^4x'\right]\cdots,$$

where $\widehat{\mathcal{H}}_{\mathrm{int}\,[x']}$ is the operator corresponding to the interaction Hamiltonian density $\mathcal{H}_{\mathrm{int}\,[x']}$.

The supplementary condition (24) will be converted to

$$\begin{aligned}\Big[\partial_\rho W^{i\rho}(x;\sigma) - \int_\sigma d\Sigma_{\rho(x'')}\big[&\partial''^{\mu} D(x-x'')\cdot g f^{ijk} W^j_{\mu}{}_{(x'')} W^{k\rho}{}_{(x'')} \\ &- D(x-x'')\cdot g f^{ijk} G^{j\rho\nu}{}_{(x'')} W^k_{\nu}{}_{(x'')} \\ &- D(x-x'')\cdot g f^{icb}[\partial'^{\rho}\bar{c}^c{}_{(x'')}]c^b{}_{(x'')} \\ &+ D(x-x'')\cdot ig\bar{\psi}_{(x'')}\gamma^\rho\lambda^i\psi_{(x'')}\big]\Big]^{(+)}|\Psi >= 0,\end{aligned} \tag{33}$$

where the symbol (+) means the positive-frequency part, and $|\Psi >$ is a state vector in the Heisenberg picture.

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